The interplay between high energy physics and cosmology: an example

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Cosmology and high energy physics are two closely connected areas. In this lecture I present an example of their rich interplay.

1. Introduction

The cosmology of the early universe is a fast progressing area of physics. This remarkable progress has been mainly achieved by employing high energy physics models.

Most aspects of high energy physics beyond the standard model can only be tested by going to very high energies, which are by far greater than those accessible by present, or even future, terrestrial accelerators. The rich interplay between particle physics and cosmology has offered a promising approach to *experimentally* test new theories of fundamental forces.

Thus, high energy physics models give us the means to formulate scenarios of the evolution of the early universe, while by confronting the predictions of these scenarios against cosmological data, one can constrain the parameters, or even falsify theories of fundamental forces.

An example of the interplay between cosmology and high energy physics is the aim of my talk.

2. The plot: Topological Defects and CMB data

2.1. Topological Defects

Many particle physics models of matter admit solutions which correspond to a class of topological defects. Under the hypothesis that we understand properly, both unification of forces, as well as big bang cosmology, we expect that topological defects could have formed naturally during phase transitions followed by spontaneously broken symmetries, in the early stages of the evolution of the universe. Among the various types of topological defects, some lead to disastrous consequences for cosmology and thus, they are undesired, while some others may play a useful rôle.

Spontaneous symmetry breaking is an old idea, described within the particle physics context in terms of the Higgs field. The Symmetry is called Spontaneously Broken (SSB) if the ground state is not invariant under the full symmetry of the Lagrangian density. Thus, the vacuum expectation value of the Higgs field is nonzero. In quantum field theories, broken symmetries are restored at high enough temperatures.

In three spatial dimensions, four different kinds of topological defects can arise. The criterion for their formation during a SSB phase transition, as well as the determination of their type, both depend on the topology of the vacuum manifold $\mathcal{M}$. The properties of $\mathcal{M}$ are usually described by the $n^{\rm th}$ homotopy group $\pi_n(\mathcal{M})$. If $\mathcal{M}$ has disconnected components, or equivalently if $\pi_0(\mathcal{M}) \neq I$, then two-dimensional defects, called *domain walls*, form. The spacetime dimension d of the defects is given in terms of the order of the nontrivial homotopy group by $d = 4 - 1 - n$. If $\mathcal{M}$ is not simply connected, in other words if $\mathcal{M}$ contains loops which cannot be continuously shrunk into a point, then *cosmic strings* form. A necessary, but not sufficient, condition for the existence of stable strings is that the fundamental group $\pi_1(\mathcal{M})$ of $\mathcal{M}$, is nontrivial, or $\mathcal{M}$ is multiply connected. Cosmic strings are line-like defects, $d = 2$. If $\mathcal{M}$ contains unshrinkable surfaces, then *monopoles* form. Finally, if $\mathcal{M}$ contains non-

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contractible three-spheres then event-like defects, *textures*, form for which $n = 3,\ d = 0$.

Depending on whether the symmetry is local (gauged) or global (rigid), topological defects are called local or global. The energy of local defects is strongly confined, while the gradient energy of global defects is spread out over the causal horizon at defect formation.

2.2. Cosmic Microwave Background Temperature Anisotropies

The Cosmic Microwave Background (CMB) temperature anisotropies provide a powerful test for theoretical models aiming at describing the early universe. The characteristics of the CMB anisotropy multipole moments, and more precisely the position and amplitude of the acoustic peaks, as well as the statistical properties of the CMB anisotropies, can be used to discriminate among theoretical models, as well as to constraint the parameters space. CMB anisotropies are characterized by their angular power spectrum C_ℓ, which is the average value of the square of the coefficients of a spherical harmonic decomposition of the measured CMB pattern.

The predictions of the defects models regarding the characteristics of the CMB spectrum are:

• Global $\mathcal{O}(4)$ textures predict the position of the first acoustic peak at $\ell \simeq 350$ with an amplitude ~ 1.5 times higher than the Sachs-Wolfe plateau [1].

• Global $\mathcal{O}(N)$ textures in the large N limit lead to a quite flat spectrum, with a slow decay after $\ell \sim 100$ [2]. Similar are the predictions of other global $\mathcal{O}(N)$ defects [3,4].

• Local cosmic strings predictions are not very well established and range from an almost flat spectrum [5] to a single wide bump at $\ell \sim 500$ [6] with extremely rapidly decaying tail.

The position and amplitude of the acoustic peaks, as found by the CMB measurements [7, 8,9,10], are in disagreement with the predictions of topological defects models.

In addition, topological defects predict nongaussian statistics of the CMB anisotropies. One could address the question whether the inflaton field can also give some nongaussian signatures. It is often assumed that the initial state of the perturbations of the inflaton field is the vacuum. In the absence of a theoretical justification for this assumption, one may relax it. The simplest way to generalise the vacuum initial state, which contains no privileged scale, is to consider [11] an initial state with a built-in characteristic scale. In a band localized around the preferred scale, the state contains a number of quanta, whereas it is still the vacuum elsewhere. A robust prediction of such a model is the nongaussian character of the induced perturbations. For models with a preferred scale, the three point (and any higher-order odd-point) correlation function vanishes, whereas the four-point (and any higher-order even-point) correlation function does not satisfy Gaussian statistics [11]. Studying such a model in the context of single-field inflation, we have shown [11,12] that the nongaussian signature is much smaller than the cosmic variance, thus undetectable. We have thus concluded that Gaussian statistics is a robust prediction of single-field inflation. From the experimental point of view, nongaussianity is strongly constrained from the WMAP measurements [13].

In conclusion, CMB measurements rule out pure topological defects models as the origin of initial density perturbations; inflation wins over topological defects. This leads to a crucial set of questions concerning high energy physics. Namely, are topological defects, and more precisely cosmic strings, allowed at all? We are basically interested in cosmic strings, since we consider gauge theories (domain walls and monopoles are dangerous, while textures are uninteresting [14]). How generic is cosmic strings formation? Which are the consequences for fundamental theories? In what follows, we address these questions.

It is conceivable to consider a mixed perturbation model, in which the primordial fluctuations are induced by an inflaton field with a nonnegligible cosmic strings contribution. We have considered [15] a model in which a network of cosmic strings evolved independently of any preexisting fluctuation background, generated by a standard cold dark matter with a nonzero cosmological constant inflationary phase. Restricting our attention to the angular spectrum, we can

remain in the linear regime. Thus,

$$C_\ell = \alpha C_\ell^{\mathrm{I}} + (1-\alpha) C_\ell^{\mathrm{S}} \ , \tag{1}$$

where C_ℓ^{I} and C_ℓ^{S} denote the (COBE normalized) Legendre coefficients due to adiabatic inflation fluctuations and those stemming from the string network respectively. The coefficient α in Eq. (1) is a free parameter giving the relative amplitude for the two contributions. One has to compare the C_ℓ, given by Eq. (1), with data obtained from CMB measurements. Already MAXIMA [7], BOOMERanG [8] and DASI [9] experiments imposed [15] an upper limit on the cosmic strings contribution to the CMB, which is $\lesssim 18\%$. Clearly, the limit set by the Wilkinson Microwave Anisotropy Probe (WMAP) measurements [10] should be stronger. A recent Bayesian analysis in a three dimensional parameter space [16] has shown that a cosmic strings contribution to the primordial fluctuations higher than 9% is excluded up to 99% confidence level.

3. The plot thickens: Genericity of cosmic strings formation in SUSY GUTs

The natural question one has to address is how generic cosmic strings formation is. Clearly the answer to this question depends on the framework we are placed in. This issue has been studied in detail within Supersymmetric Grand Unified Theories (SUSY GUTs) in Ref. [17].

Grand Unified Theories imply that our universe has undergone a series of phase transitions associated with the SSB of the GUT gauge group $\mathrm{G_{GUT}}$ down to the standard model gauge group $\mathrm{G_{SM}} = \mathrm{SU(3)_C} \times \mathrm{SU(2)_L} \times \mathrm{U(1)_Y}$ at $M_{\mathrm{GUT}} \sim 3 \times 10^{16}$ GeV. There might be one, more than one, or no intermediate symmetry group between $\mathrm{G_{GUT}}$ and $\mathrm{G_{SM}}$. As a cosmological consequence of these SSB patterns one obtains the formation of topological defects via the Kibble mechanism [18].

Spontaneous symmetry breaking schemes which lead to the formation of monopoles or domain walls are ruled out since they are incompatible with our universe, unless an inflationary era took place after their formation. We cannot constrain SSB schemes with texture formation, since this class of defects cannot play a significant rôle in cosmology [14].

The particle physics Standard Model (SM) has been tested to a very high precision, however experimental data, and in particular evidence of neutrino masses [19,20,21], show that one should go beyond this model. An extension of the SM gauge group is realised in the framework of Supersymmetry (SUSY), which is at present the only viable theory for solving the gauge hierarchy problem. In addition, within SUSY GUTs the gauge coupling constants of the strong, weak and electromagnetic interactions meet at a single point $M_{\mathrm{GUT}} \simeq (2-3) \times 10^{16}$ GeV. Finally, SUSY GUTs can provide the scalar field to play the rôle of an inflaton field, explain the baryon asymmetry of the universe, and provide a candidate (the lightest superparticle) for cold dark matter.

In Ref. [17], we have considered all possible SSB schemes from a large gauge group $\mathrm{G_{GUT}}$ down to $\mathrm{G_{SM}} \times Z_2$, in the context of SUSY GUTs. Z_2 is a sub-group of the $\mathrm{U(1)_{B-L}}$ gauge symmetry and it plays the rôle of R-parity. The requirement of an unbroken R-parity down to low energies guarantees proton stability. We have limited the choice of $\mathrm{G_{GUT}}$ to simple gauge groups which contain $\mathrm{G_{SM}}$, have a complex representation, are anomaly free, and whose rank is not higher than 8; the main conclusions remain qualitatively unaffected for groups of higher rank. We have studied the homotopy group of the vacuum manifold to find the type of defects formed, if any.

We have considered an exhaustive list of possible embeddings of $\mathrm{G_{SM}}$ in $\mathrm{G_{GUT}}$ and we have examined whether defects are formed during the SSB patterns and of which kind they are. To get rid of the undesired defects, basically monopoles, we have employed an era of standard hybrid inflation after their formation. Moreover, we have considered a mechanism of baryogenesis via leptogenesis, which can be thermal or nonthermal one. In the case of nonthermal leptogenesis, $\mathrm{U(1)_{(B-L)}}$ is a sub-group of the GUT gauge group, $\mathrm{G_{GUT}}$, and B-L is broken at the end or after inflation. In the case of thermal leptogenesis, B-L is broken independently of inflation. If leptogenesis is thermal and B-L is broken before the inflationary era, then one should check whether the temperature at which B-L is broken, which will define

the mass of the right-handed neutrinos, is smaller than the reheating temperature which should be lower than the limit imposed by the gravitino.

We have then asked how generic is cosmic strings formation after hybrid inflation, within these schemes. Only if we relax the requirement that the gauged $\mathrm{B}-\mathrm{L}$ symmetry is broken at the end of inflation, there are a few ($\sim 2\%$) SSB schemes without cosmic strings formation, otherwise cosmic strings formation is generic. The number of SSB schemes with no cosmic strings formation increases ($\sim 15\%$), if we accept SSB schemes with broken R-parity, but then it remains an open question how to stabilise the proton.

We have found [17] that cosmic strings formation is sometimes accompanied by the formation of embedded strings which however are topologically, and in general also dynamically, unstable [22].

4. The plot unfolds: Supersymmetric Hybrid Inflation

One of the main questions within our list is indeed answered. Cosmic strings are generically formed in the framework of SUSY GUTs. However, we stated earlier that strong constraints are placed in their contribution to the CMB power spectrum. Thus, the obvious question to address at this point, is whether we can constrain the parameter's space of the models with strings at the end of the last inflationary era, so that their contribution to the CMB is within the allowed window. Answering to this question will allow us, at least partially, to find the class of natural inflationary model, if any. These issues have been addressed in Refs. [23,24,25].

The inflationary paradigm offers the most appealing approach for describing the early stages of the evolution of our universe. Inflation essentially consists of a phase of accelerated expansion which took place at a very high energy scale. Inflation requires the existence of a slowly rolling scalar field, while inflation will cease whenever slow-roll conditions are violated. Inflation comes to complete the standard Big Bang model and offers an explanation for the initial density fluctuations leading to the observed structure formation and the measured anisotropies of the CMB. However, inflation is faced with two questions, namely how generic is the onset of inflation and which is a natural model of inflation. We have found [26] that the onset of inflation requires some special initial conditions, which however may be the likely outcome of quantum events occurred before the inflationary era [27]. To find a natural model of inflation, consistent with high energy physics models and cosmological data seems to be less trivial.

To describe the early evolution of our universe, at energies below Planck scale, one should consider an effective N=1 Supergravity (SUGRA). This implies that inflationary models should be constructed in the framework of SUGRA, since the inflationary scale is $V^{1/4} \lesssim 4 \times 10^{16}$ GeV. However, is is difficult to implement slow-roll inflation within SUGRA. More precisely, the positive false vacuum of the inflaton field breaks spontaneously global supersymmetry; it gets restored after the end of the inflationary era, when the field rolls to the true vacuum. In SUGRA the SUSY breaking is transmitted to all fields by gravity, thus any scalar field gets a soft mass given by

$$m_{\rm soft}^2 \sim 8\pi V/M_{\rm Pl}^2 \sim H^2 , \tag{2}$$

where H is the expansion rate during inflation, and $M_{\rm Pl}$ denotes the reduced Planck mass. One has to use fine-tuning to avoid such a large soft mass for the scalar field which plays the rôle of the inflaton. This is known as the problem of "Hubble-induced mass".

In a supersymmetric theory, the tree-level potential is the sum of an F-term and a D-term, which have different properties. In all proposed inflationary models one of these two terms is the dominant one. It was shown [28], that the "Hubble-induced mass" problem comes from F-term interactions and it may be avoided if we consider the vacuum energy as being dominated by nonzero D-terms of some superfields. Inflationary models where the potential is dominated by nonvanishing D-terms emerge naturally in theories with either an anomalous or an nonanomalous gauge U(1) symmetry which incorporates a Fayet-Iliopoulos term. In D-term inflation the masses of the scalar fields depend on their gauge charges.

More precisely, in D-term inflation, the inflaton field is a singlet under gauge symmetry, thus the curvature of the inflaton potential is small. Since in addition, D-term inflation can easily be implemented in string theory, this class of models gained a lot of interest.

4.1. F-term Inflation

F-term inflation can be accommodated in a SSB scheme, where a GUT gauge group is broken down to the standard model gauge group at an energy scale $M_{\rm GUT}$ according to

$$G_{\rm GUT} \xrightarrow{M_{\rm GUT}} H_1 \xrightarrow[\Phi_+\Phi_-]{M_{\rm infl}} H_2 \longrightarrow G_{\rm SM} \ , \qquad (3)$$

where Φ_+, Φ_- is a pair of GUT Higgs superfields in nontrivial complex conjugate representations, which lower the rank of the group by one unit when acquiring nonzero VEV. The inflationary phase takes place at the beginning of the SSB $H_1 \xrightarrow{M_{\rm infl}} H_2$. F-term inflation can be chosen as the hybrid inflationary model introduced in the SSB schemes studied in the previous section. Thus, generically, one expects the formation of cosmic strings (accompanied sometimes by embedded strings), at the end of the inflationary era.

F-term inflation is based on the supersymmetric renormalisable superpotential

$$W_{\rm infl}^{\rm F} = \kappa S(\Phi_+\Phi_- - M^2) \ , \qquad (4)$$

where S, Φ_+, Φ_- are chiral superfields, and κ, M are two constants. The scalar potential reads

$$V(\phi_+,\phi_-,S) = |F_{\Phi_+}|^2 + |F_{\Phi_-}|^2 + |F_S|^2 + \frac{1}{2}\sum_a g_a^2 D_a^2 \ . \qquad (5)$$

The F-term is such that $F_{\Phi_i} \equiv |\partial W/\partial\Phi_i|_{\theta=0}$, where we take the scalar component of the superfields once we differentiate with respect to $\Phi_i = \Phi_+, \Phi_-, S$. The D-term is

$$D_a = \bar{\phi}_i \, (T_a)^i{}_j \, \phi^j + \xi_a \ , \qquad (6)$$

with a the label of the gauge group generators T_a, g_a the gauge coupling, and ξ_a the Fayet-Iliopoulos term. By definition, in the F-term inflation the real constant ξ_a is zero; it can only be nonzero if T_a generates a U(1) group.

In the context of F-term hybrid inflation, the F-terms give rise to the inflationary potential energy density, while the D-terms are flat along the inflationary trajectory, thus one may neglect them during inflation.

The potential has one valley of local minima, $V = \kappa^2 M^4$, for $S > M$ with $\phi_+ = \phi_- = 0$, and one global supersymmetric minimum, $V = 0$, at $S = 0$ and $\phi_+ = \phi_- = M$. Imposing initially $S \gg M$, the fields quickly settle down the valley of local minima. Since in the slow roll inflationary valley, the ground state of the scalar potential is nonzero, SUSY is broken. In the tree level, along the inflationary valley the potential being constant, it is perfectly flat. A slope along the potential can be generated by including the one-loop radiative corrections. Thus, the scalar potential gets a little tilt which helps the inflaton field S to slowly roll down the valley of minima. The one-loop radiative corrections to the scalar potential along the inflationary valley, lead to an effective potential [29,30,31,23,24,25]

$$\begin{aligned} V_{\rm eff}^{\rm F}(|S|) = \ & \kappa^2 M^4 \Big\{ 1 + \frac{\kappa^2 \mathcal{N}}{32\pi^2} \Big[2 \ln \frac{|S|^2\kappa^2}{\Lambda^2} \\ & + \Big(\frac{|S|^2}{M^2} + 1\Big)^2 \ln \Big(1 + \frac{M^2}{|S|^2}\Big) \Big) \\ & + \Big(\frac{|S|^2}{M^2} - 1\Big)^2 \ln \Big(1 - \frac{M^2}{|S|^2}\Big) \Big) \Big] \Big\} \ ; \end{aligned} \qquad (7)$$

Λ is a renormalisation scale and $\mathcal{N}$ stands for the dimensionality of the representation to which the complex scalar components ϕ_+, ϕ_- of the chiral superfields Φ_+, Φ_- belong.

Considering only large angular scales, one can study the contributions to the CMB temperature anisotropies by analytical methods. In Refs. [23,24], we have calculated explicitly the Sachs-Wolfe effect. The quadrupole anisotropy has one contribution coming from the inflaton field, calculated using Eq. (7), and one contribution coming from the cosmic strings network, given by numerical simulations [32]. Fixing the number of e-foldings to 60, then for a given gauge group, the inflaton and cosmic strings contribution to the CMB depend on the superpotential coupling κ, or equivalently on the symmetry breaking scale M associated with the inflaton

mass scale, which coincides with the string mass scale. The total quadrupole anisotropy has to be normalised to the COBE data. In Refs. [23,24] we have found that the cosmic strings contribution is consistent with the CMB measurements, provided

$$M \lesssim 2 \times 10^{15}\mathrm{GeV} \quad \Leftrightarrow \quad \kappa \lesssim 7 \times 10^{-7} \ . \tag{8}$$

Strictly speaking the above condition was found in the context of SO(10) gauge group, but the conditions imposed in the context of other gauge groups are of the same order of magnitude since M is a slowly varying function of the dimensionality $\mathcal{N}$ of the representations to which the scalar components of the chiral Higgs superfields belong.

The superpotential coupling κ is also subject to the gravitino constraint which imposes an upper limit to the reheating temperature, to avoid gravitino overproduction. Within the framework of SUSY GUTs and assuming a see-saw mechanism to give rise to massive neutrinos, the inflaton field will decay during reheating into pairs of right-handed neutrinos. This constraint on the reheating temperature can be converted to a constraint on the parameter κ. The gravitino constraint on κ reads [23,24] $\kappa \lesssim 8 \times 10^{-3}$, which is a weaker constraint.

Concluding, F-term inflation leads generically to cosmic strings formation at the end of the inflationary era. The cosmic strings formed are of the GUT scale. This class of models can be compatible with CMB measurements, provided the superpotential coupling is smaller than 10^{-6}. This tuning on the free parameter κ can be softened if one allows for the curvaton mechanism. According to the curvaton mechanism [33,34], another scalar field, called the curvaton, could generate the initial density perturbations whereas the inflaton field is only responsible for the dynamics of the universe. The curvaton is a scalar field, that is subdominant during the inflationary era as well as at the beginning of the radiation dominated era which follows the inflationary phase. There is no correlation between the primordial fluctuations of the inflaton and curvaton fields. Clearly, within supersymmetric theories such scalar fields are expected to exist. In this case, the coupling κ is only constrained by the gravitino limit. More precisely, assuming the existence of a curvaton field, there is an additional contribution to the temperature anisotropies. The WMAP CMB measurements impose [23,24] the following limit on the initial value of the curvaton field

$$\psi_{\mathrm{init}} \lesssim 5 \times 10^{13} \left(\frac{\kappa}{10^{-2}}\right) \mathrm{GeV} \ , \tag{9}$$

provided the parameter κ is in the range $[10^{-6},\ 1]$.

4.2. D-term Inflation

D-term inflation is derived in SUSY from the superpotential

$$W^{\mathrm{D}}_{\mathrm{infl}} = \lambda S \Phi_+ \Phi_- \ , \tag{10}$$

where S, Φ_-, Φ_+ are chiral superfields and λ is the superpotential coupling. In D-term inflation, as opposed to F-term inflation, the inflaton mass acquires values of the order of Planck mass, and therefore, the correct analysis must be done in the framework of SUGRA.

For minimal SUGRA, the effective scalar potential reads [23,24,25]

$$\begin{aligned} V^{\mathrm{D-SUGRA}}_{\mathrm{eff}} &= \frac{g^2\xi^2}{2}\Big\{1 + \frac{g^2}{16\pi^2} \\ &\times \Big[2\ln\frac{|S|^2\lambda^2}{\Lambda^2} e^{\frac{|S|^2}{M^2_{\mathrm{Pl}}}} \\ &+ \Big(\frac{\lambda^2|S|^2}{g^2\xi} e^{\frac{|S|^2}{M^2_{\mathrm{Pl}}}} + 1\Big)^2 \ln\Big(1 + \frac{g^2\xi}{\lambda^2|S|^2} e^{-\frac{|S|^2}{M^2_{\mathrm{Pl}}}}\Big) \\ &+ \Big(\frac{\lambda^2|S|^2}{g^2\xi} e^{\frac{|S|^2}{M^2_{\mathrm{Pl}}}} - 1\Big)^2 \ln\Big(1 - \frac{g^2\xi}{\lambda^2|S|^2} e^{-\frac{|S|^2}{M^2_{\mathrm{Pl}}}}\Big)\Big]\Big\} \end{aligned} \tag{11}$$

D-term inflation requires a SSB pattern,

$$G_{\mathrm{GUT}} \times U(1) \xrightarrow{M_{\mathrm{GUT}}} H \times U(1) \xrightarrow[\Phi_+\Phi_-]{M_{\mathrm{infl}}} H \rightarrow G_{\mathrm{SM}} \ . \tag{12}$$

Clearly, the symmetry breaking at the end of the inflationary phase implies that cosmic strings are always formed at the end of D-term hybrid inflation. This statement, thought to cause a problem for D-term inflation, since it was claimed [35] that in this class of models, the cosmic strings contribution to the CMB measurements is constant and dominant. In the literature, one can find a number of approaches to avoid cosmic strings formation in the context of D-term inflation. For

example, one can add a nonrenorlisable term in the potential [36], or add an additional discrete symmetry [37], or consider GUT models based on nonsimple groups [38], or introduce a new pair of charged superfields [39] so that cosmic strings formation is avoided within D-term inflation.

In Refs. [23,24,25], we have properly addressed the question of cosmic strings contribution to the CMB data and we have found that standard D-term inflation can be compatible with measurements; the cosmic strings contribution to the CMB is actually model-dependent. Our most important finding was that cosmic strings contribution is not constant, nor is it always dominant.

More precisely, we have found [23,24,25] that $g \gtrsim 2 \times 10^{-2}$ is incompatible with the allowed cosmic strings contribution to the WMAP measurements. For $g \lesssim 2 \times 10^{-2}$, the constraint on the superpotential coupling λ reads $\lambda \lesssim 3 \times 10^{-5}$. SUGRA corrections impose in addition a lower limit to λ. The constraints induced on the couplings by the CMB measurements can be expressed as a single constraint on the Fayet-Iliopoulos term ξ, namely $\sqrt{\xi} \lesssim 2 \times 10^{15}$ GeV.

Concluding, standard D-term inflation always leads to cosmic strings formation at the end of the inflationary era. The cosmic strings formed are of the GUT scale. This class of models is still compatible with CMB measurements, provided the couplings are small enough. As in the previous class of models, the fine tuning on the couplings can be softened provided one considers the curvaton mechanism. In this case, the imposed CMB constraint on the initial value of the curvaton field reads [25]

$$\psi_{\rm init} \lesssim 3 \times 10^{14} \left(\frac{g}{10^{-2}}\right) \ {\rm GeV}, \tag{13}$$

for $\lambda \in [10^{-1}, 10^{-4}]$.

5. Conclusions

High energy physics models are used to formulate scenarios of the evolution of the early universe. Comparing the predictions of these cosmological scenarios against observational and experimental data we can in return test the original high energy physics models.

Cosmic strings are generically formed in almost all SSB schemes from a large gauge group down to the standard model. To be consistent with CMB measurements, we can place limits on the free parameters (masses and couplings). Thus, F- as well as D-term inflationary models are found to be consistent with the data, provided the couplings are small enough, unless we want to employ the curvaton mechanism.

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